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Perturbative QCD study of B_s decays to a pseudoscalar meson and a tensor meson

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ABSTRACT

We study two-body hadronic $B_s \to PT$ decays, with $P(T)$ being a light pseudoscalar (tensor) meson, in the perturbative QCD approach. The CP-averaged branching ratios and the direct CP asymmetries of the $\Delta S = 0$ modes are predicted, where ΔS is the difference between the strange numbers of final and initial states. We also define and calculate experimental observables for the $\Delta S = 1$ modes under the B_s^0–$\bar{B}_s^0$ mixing, including CP averaged branching ratios, time-integrated CP asymmetries, and the CP observables C_f, D_f and S_f. Results are compared to the $B_s \to PV$ ones in the literature, and to the $B \to PT$ ones, which indicate considerable U-spin symmetry breaking. Our work provides theoretical predictions for the $B_s \to PT$ decays for the first time, some of which will be potentially measurable at future experiments.

Two-body hadronic B meson decays have attracted a lot of attention, because of their importance for studies of CP violation, CKM angle determination, and both weak and strong dynamics. The two B factories have measured hadronic B decays into light tensor (T) mesons recently [1–3], which were also intensively investigated in several theoretical methods, such as the naive factorization hypothesis [4–6], the perturbative QCD (PQCD) approach [7], and the QCD factorization approach [8]. With much higher production efficiency of B_s mesons at the LHCb than at the B factories, many data for two-body hadronic B_s decays have been published [9,10], but no decays into tensor mesons were observed so far.

The B_s decays into tensor mesons have not been analyzed theoretically either, to our knowledge. The naive factorization hypothesis does not apply to modes involving only the annihilation amplitudes and only the amplitudes with tensor mesons being emitted from the weak vertex. Besides, branching ratios for color-suppressed decays estimated in the naive factorization are usually too small. As for the QCD factorization [11], owning to the lack of data for $B_s \to PT$ branching ratios, P being a light pseudoscalar meson, the penguin-annihilation parameters cannot be determined through global fits. If the parameters associated with the $B_s \to PT$ modes were approximated by the $B_s \to PV$ ones [8], large theoretical uncertainties would be introduced. Both the annihilation amplitudes and the nonfactorizable tensor-emission amplitudes are calculable in the PQCD approach without inputs of free parameters. Encouraged by successful applications of the PQCD approach to many two-body hadronic B meson decays [12,14,7,13], in this Letter we will make predictions for the $B_s \to PT$ branching ratios and CP-violation observables, which can provide useful hints to relevant experiments.

The effective electroweak Hamiltonian relevant to the $B_s \to PT$ decays is written as

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}}\left[\sum_{i=1}^{2} V_{ub}^* V_{uD} C_i(\mu) O_i^u(\mu) - \sum_{j=3}^{10} V_{tb}^* V_{tD} C_j(\mu) O_j^u(\mu)\right], \tag{1}$$

where V's are the CKM matrix elements with D denoting a down-type quark d or s, $O_{i,j}(\mu)$ are the tree and penguin four-quark operators [15], and $C_{i,j}(\mu)$ are the corresponding Wilson coefficients, which evolve from the W boson mass down to the renormalization scale μ. In the PQCD approach a hadronic transition matrix element of a four-quark operator is further factorized into two pieces [16]: the kernel with hard gluon exchanges characterized

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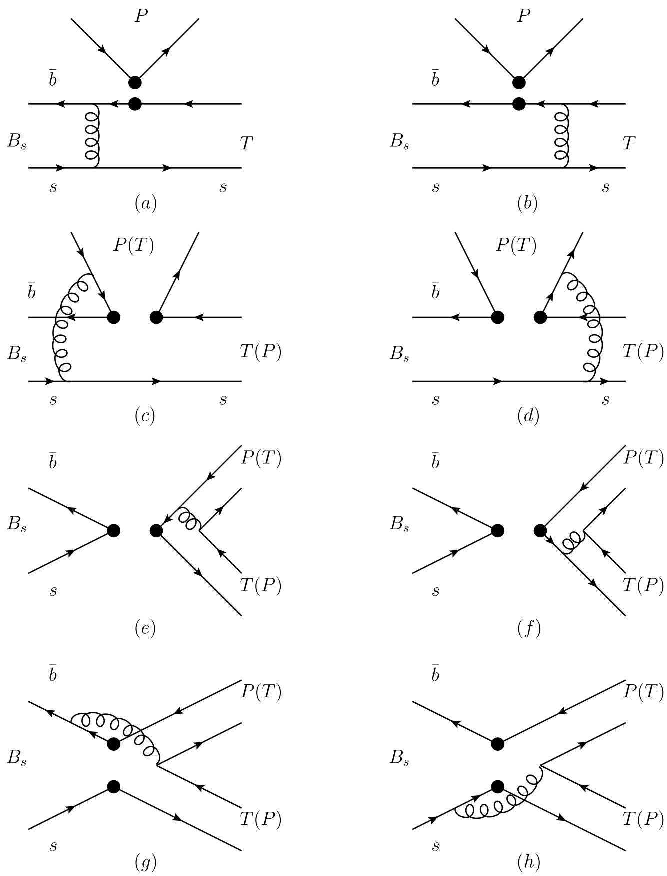

Fig. 1. Leading-order diagrams for $B_s \to PT$ decays.

by the b quark mass and the nonperturbative hadron wave functions characterized by the QCD scale $\Lambda_{\rm QCD}$.

The leading-order diagrams contributing to the $B_s \to PT$ decays are displayed in Fig. 1, where (a) and (b) are factorizable emission-type diagrams, (c) and (d) are nonfactorizable emission-type diagrams, (e) and (f) are factorizable annihilation-type diagrams, and (g) and (h) are nonfactorizable annihilation-type diagrams. As indicated in Fig. 1, the factorizable tensor-emission amplitudes do not exist, since a tensor meson cannot be produced via a V or A current. The PQCD results for the $B \to PT$ (without B_s) decays [7] are basically in agreement with the experimental data [17,18] and those from the QCD factorization [8]. The extension of the PQCD formalism to the $B_s \to PT$ decays is straightforward because of the similarity between B and B_s decays in $SU(3)$ symmetry: the factorization formula for every diagram can be obtained by substituting the quantities in the $B_s \to PT$ decays for the corresponding ones in the $B \to PT$ decays [7]. The confrontation of the $B \to PT$ calculations to the data has restricted the parameters involved in the P and T meson wave functions to some extent. In this work we will adopt the B_s meson wave function in [14], and the P and T meson wave functions in [7].

A neutral meson and its charge conjugate partner, including the $K^0-\bar K^0$, $D^0-\bar D^0$, $B^0-\bar B^0$, and $B_s^0-\bar B_s^0$ systems, mix through the weak interaction. The $B_s^0-\bar B_s^0$ mixing is the strongest, since the mass difference ΔM between the mass eigenstates is much larger than the decay width Γ of the B_s meson. The frequent oscillation between the B_s^0 and $\bar B_s^0$ mesons due to the strong mixing has rendered difficult measurements of B_s decay observables at the B factories, such as measurements of time-dependent CP-violation parameters. However, these measurements become feasible in LHCb experiments, because of the time dilation caused by energetic B_s mesons. The mass eigenstates of the B_s mesons are superpositions of the flavor eigenstates,

$$|B_{sL,H}\rangle = p\big|B_s^0\rangle \pm q\big|\bar B_s^0\rangle, \tag{2}$$

where p and q are complex coefficients. We neglect the difference between the mass eigenstates and the CP eigenstates, and assume that $B_{sL(H)}$ is CP even (odd) as suggested in [19]. The time-dependent $B_s \to PT$ differential branching ratios are then expressed as [20]

$$\begin{aligned}
\frac{d}{dt}Br\big(B_s^0(t)\to f\big) &= \Phi(B_s\to f)e^{-\Gamma t}|A_f|^2\frac{1+|\lambda_f|^2}{2}\\
&\quad\times\left[\cosh\left(\frac{\Delta\Gamma}{2}t\right)+\cos(\Delta M t)C_f\right.\\
&\quad\left.-\sin(\Delta M t)S_f-\sinh\left(\frac{\Delta\Gamma}{2}t\right)D_f\right],\\
\frac{d}{dt}Br\big(\bar B_s^0(t)\to f\big) &= \Phi(B_s\to f)e^{-\Gamma t}\left|\frac{p}{q}\right|^2|A_f|^2\frac{1+|\lambda_f|^2}{2}\\
&\quad\times\left[\cosh\left(\frac{\Delta\Gamma}{2}t\right)-\cos(\Delta M t)C_f\right.\\
&\quad\left.+\sin(\Delta M t)S_f-\sinh\left(\frac{\Delta\Gamma}{2}t\right)D_f\right],\\
\frac{d}{dt}Br\big(\bar B_s^0(t)\to \bar f\big) &= \Phi(B_s\to f)e^{-\Gamma t}|\bar A_{\bar f}|^2\frac{1+|\bar\lambda_{\bar f}|^2}{2}\\
&\quad\times\left[\cosh\left(\frac{\Delta\Gamma}{2}t\right)+\cos(\Delta M t)C_{\bar f}\right.\\
&\quad\left.-\sin(\Delta M t)S_{\bar f}-\sinh\left(\frac{\Delta\Gamma}{2}t\right)D_{\bar f}\right],\\
\frac{d}{dt}Br\big(B_s^0(t)\to \bar f\big) &= \Phi(B_s\to f)e^{-\Gamma t}\left|\frac{q}{p}\right|^2|\bar A_{\bar f}|^2\frac{1+|\bar\lambda_{\bar f}|^2}{2}\\
&\quad\times\left[\cosh\left(\frac{\Delta\Gamma}{2}t\right)-\cos(\Delta M t)C_{\bar f}\right.\\
&\quad\left.+\sin(\Delta M t)S_{\bar f}-\sinh\left(\frac{\Delta\Gamma}{2}t\right)D_{\bar f}\right],
\end{aligned}\tag{3}$$

with the mass difference $\Delta M = (116.4 \pm 0.5)\times 10^{-10}$ MeV, the decay width difference $\Delta\Gamma = (0.100 \pm 0.013)\times 10^{12}\ \text{s}^{-1}$ [17], $\Phi(B_s \to f)$ being the phase space of the corresponding mode, and A_f ($\bar A_{\bar f}$) being the $B_s^0 \to f$ ($\bar B_s^0 \to \bar f$) decay amplitude. We have employed the definitions of the amplitude ratios λ_f and $\bar\lambda_{\bar f}$, and the CP asymmetry observables $C_{f,\bar f}$, $D_{f,\bar f}$ and $S_{f,\bar f}$ used in [20].

Since the oscillation period is much shorter than the lifetime of the B_s meson, Eq. (3) can be integrated over t, and lead to the time-integrated branching ratios

$$\begin{aligned}
&Br\big(B_s^0(\infty)\to f\big)\\
&\quad= \Phi(B_s\to f)|A_f|^2\frac{1+|\lambda_f|^2}{2}\left[\frac{\Gamma-D_f\frac{\Delta\Gamma}{2}}{\Gamma^2}+\frac{C_f\Gamma+S_f\Delta M}{\Gamma^2+\Delta M^2}\right],\\
&Br\big(\bar B_s^0(\infty)\to f\big)\\
&\quad= \Phi(B_s\to f)|A_f|^2\frac{1+|\lambda_f|^2}{2}\left[\frac{\Gamma-D_f\frac{\Delta\Gamma}{2}}{\Gamma^2}-\frac{C_f\Gamma+S_f\Delta M}{\Gamma^2+\Delta M^2}\right],\\
&Br\big(\bar B_s^0(\infty)\to \bar f\big)\\
&\quad= \Phi(B_s\to f)|\bar A_{\bar f}|^2\frac{1+|\bar\lambda_{\bar f}|^2}{2}\left[\frac{\Gamma-D_{\bar f}\frac{\Delta\Gamma}{2}}{\Gamma^2}+\frac{C_{\bar f}\Gamma+S_{\bar f}\Delta M}{\Gamma^2+\Delta M^2}\right],
\end{aligned}$$

Table 1
Branching ratios (in units of 10^{-7}) and direct CP asymmetries of the $\Delta S = 0$ $B_s^0 \to PT$ decays.

Modes	Amplitudes	Br	Direct A_{CP} (%)
$B_s^0 \to \pi^+ K_2^{*-}$	T	90^{+40+4}_{-32-6}	13^{+2+2}_{-2-2}
$B_s^0 \to \pi^0 \bar{K}_2^{*0}$	C, PA	$1.3^{+0.6+0.6}_{-0.5-0.5}$	47^{+8+9}_{-6-6}
$B_s^0 \to \bar{K}^0 a_2^0$	C, PA	$2.0^{+0.4+0.2}_{-0.3-0.3}$	38^{+7+6}_{-10-7}
$B_s^0 \to \bar{K}^0 f_2$	C, PA	$3.4^{+0.7+0.7}_{-0.6-0.7}$	-24^{+5+3}_{-6-5}
$B_s^0 \to \bar{K}^0 f_2'$	PA	$2.0^{+0.5+0.8}_{-0.4-0.6}$	$4.8^{+2.8+1.9}_{-1.7-1.4}$
$B_s^0 \to K^- a_2^+$	T, PA	$1.5^{+0.3+0.4}_{-0.2-0.3}$	39^{+8+1}_{-1-4}
$B_s^0 \to \eta \bar{K}_2^{*0}$	C, PA	$0.55^{+0.29+0.35}_{-0.19-0.27}$	77^{+13+5}_{-12-2}
$B_s^0 \to \eta' \bar{K}_2^{*0}$	C, PT	$3.5^{+1.2+1.4}_{-1.0-1.2}$	-30^{+2+7}_{-1-6}

$$Br\big(B_s^0(\infty) \to \bar{f}\big)$$
$$= \Phi(B_s \to f)|\bar{A}_{\bar{f}}|^2 \frac{1+|\bar{\lambda}_{\bar{f}}|^2}{2}\left[\frac{\Gamma - D_{\bar{f}}\frac{\Delta\Gamma}{2}}{\Gamma^2} - \frac{C_{\bar{f}}\Gamma + S_{\bar{f}}\Delta M}{\Gamma^2+\Delta M^2}\right]. \quad (4)$$

The terms proportional to $(\Delta\Gamma/\Gamma)^2 \approx 0.006$ have been dropped, and the approximation $|p/q|^2 = 1$ has been made in the above expressions. If it happens that the B_{sL} state is CP odd while B_{sH} is CP even, the substitutions $\Delta M \to -\Delta M$ and $\Delta\Gamma \to -\Delta\Gamma$, or equivalently, $D_{f,\bar{f}} \to -D_{f,\bar{f}}$ and $S_{f,\bar{f}} \to -S_{f,\bar{f}}$ need to be done.

For $\Delta S = 0$ modes, a B_s^0 ($\bar{B}_s^0$) meson decays to the final state f ($\bar{f}$), but not to $\bar{f}$ (f) with $f \neq \bar{f}$. In this case one can determine the initial B_s^0 or $\bar{B}_s^0$ meson through the final state even under the frequent $B_s^0 - \bar{B}_s^0$ oscillation. The ordinary definitions of CP-averaged branching ratios and direct CP asymmetries then apply directly. The predictions for the CP-averaged branching ratios and the direct CP asymmetries of these $\Delta S = 0$ modes are listed in Table 1. The dominant topological amplitudes for each decay channel are also listed, including the color-favored (T), color-suppressed (C), and annihilation-type (A) tree amplitudes, and the corresponding penguin amplitudes PT, PC, and PA. Two types of theoretical uncertainties are estimated here: the first type comes from the variation of the nonperturbative parameters in the meson wave functions (see [7,14], except that we have adopted the recent lattice QCD result for the B_s meson decay constant, 0.228(10) GeV [21]); the second type reflects the unknown next-to-leading-order QCD corrections characterized by the variations of the QCD scale $\Lambda_{\mathrm{QCD}} = (0.25 \pm 0.05)$ GeV and of the hard scales. It is observed that both types of uncertainties are roughly of the same order for most channels.

As shown in Table 1, only the $B_s^0 \to \pi^+ K_2^{*-}$ decay has a sizable branching ratio arising from the dominant amplitude T, and the branching ratios of the other modes are of order 10^{-7}. For color-suppressed modes such as $B_s^0 \to \bar{K}^0 a_2^0$, $\bar{K}^0 f_2$ and $\bar{K}^0 f_2'$, there is no significant difference between their branching ratios and those of their PV partners [14], because the factorizable emission contributions are less important. For the color-favored $B_s^0 \to K^- a_2^+$ decay, whose factorizable tensor-emission amplitude is forbidden, its branching ratio 1.50×10^{-7} is much smaller than the $B_s^0 \to K^- \rho^+$ one, 1.78×10^{-5}. Most modes in Table 1 exhibit large direct CP asymmetries caused by the interference between the tree and penguin amplitudes. The direct CP asymmetry in the $B_s^0 \to \bar{K}^0 f_2'$ decay would vanish, if f_2' was a pure $\bar{s}s$ state. After receiving a tree contribution from the mixing of the isospin-1 states, this mode gets a small CP asymmetry.

To examine whether the U-spin symmetry holds in the $B_{(s)} \to PT$ decays, we define the following ratios

$$R_{CP}\big(B_s^0 \to f\big) \equiv -\frac{A_{CP}(B_s^0 \to f)}{A_{CP}(B^0 \to Uf)},$$
$$R_{\Gamma}\big(B_s^0 \to f\big) \equiv \frac{\tau(B_s^0)}{\tau(B^0)}\frac{Br(B^0 \to Uf)}{Br(B_s^0 \to f)}, \quad (5)$$

where U stands for the U-spin transformation, $d \leftrightarrow s$. The relation between two decay modes in a U-spin pair implies that the above ratios are equal to each other in the U-spin symmetry limit [22]. Combing our predictions with the $B \to PT$ ones [7], we obtain $R_{CP}(B_s^0 \to \pi^+ K_2^{*-}) = 0.29^{+0.10}_{-0.08}$ and $R_{\Gamma}(B_s^0 \to \pi^+ K_2^{*-}) = 0.74^{+0.24}_{-0.19}$; $R_{CP}(B_s^0 \to K^- a_2^+) = 1.9^{+0.5}_{-0.5}$ and $R_{\Gamma}(B_s^0 \to K^- a_2^+) = 5.2^{+0.9}_{-0.6}$. The central values indicate that the U-spin symmetry is considerably broken in the $B_{(s)} \to PT$ decays by hadronic effects at order $(m_s - m_d)/\Lambda_{\mathrm{QCD}}$ [22], m_s (m_d) being the strange (down) quark mass. The physical U-spin conjugate processes of the other modes do not exist due to the superposition of the flavor states $\bar{q}q$ in final-state mesons.

For $\Delta S = 1$ B_s^0 ($\bar{B}_s^0$) meson decays, we first consider those modes, whose final states are CP eigenstates, i.e. $f = \bar{f}$. In this case the four equations in Eq. (3) reduce to two, and one has to measure the CP observables C_f, D_f and S_f through time-dependent branching ratios, which require a lot of data accumulation. Alternatively, we define the time-integrated CP asymmetries for these decays

$$A_{CP}\big(B_s(\infty) \to f\big) \equiv \frac{Br(\bar{B}_s^0(\infty) \to f) - Br(B_s^0(\infty) \to f)}{Br(\bar{B}_s^0(\infty) \to f) + Br(B_s^0(\infty) \to f)}$$
$$= -\frac{C_f\Gamma + S_f\Delta M}{\Gamma^2+\Delta M^2}\frac{\Gamma^2}{\Gamma - D_f\frac{\Delta\Gamma}{2}}, \quad (6)$$

and assess if there is a chance to measure it at the early stage of data accumulation.

The PQCD predictions for all the experimental observables, together with the dominant topological amplitudes and uncertainties, are shown in Table 2. It is observed that the η'-involved modes $B_s^0 \to \eta' a_2^0 (f_2, f_2')$ have branching ratio larger than those of the corresponding η-involved modes $B_s^0 \to \eta a_2^0 (f_2, f_2')$. This pattern is understood, since the dominant amplitudes require the $\bar{s}s$ constituent, which is more in η' than in η. The branching ratios of the $\Delta I = 1$ modes, like $B_s^0 \to \eta a_2^0$ and $\eta' a_2^0$, are highly suppressed, compared to those of the corresponding $\Delta I = 0$ modes, $B_s^0 \to \eta f_2$ and $\eta' f_2$. This suppression can be explained as follows. Neglecting the f_2–f_2' mixing effect, both $B_s^0 \to \eta' a_2^0$ and $\eta' f_2$ are dominated by the amplitudes PC naively. However, the minus sign in the flavor constituent $(\bar{u}u - \bar{d}d)/\sqrt{2}$ renders $PC(u)$ and $PC(d)$ cancel in the former mode, while they become constructive in the latter. The source of the discrepancy between the $B_s^0 \to \eta a_2^0$ and ηf_2 branching ratios is the same.

Contrary to the $\Delta S = 0$ decays, the tree and penguin contributions are never simultaneously sizable to form significant interferences in the $\Delta S = 1$ decays listed in Table 2, so the direct CP asymmetries C_f's are tiny. One seemingly exceptional mode is $B_s^0 \to \pi^0 f_2$, which has the tree and penguin contributions of the same order, but still a small direct CP asymmetry. A careful investigation reveals that the strong phases of the tree and penguin amplitudes are almost equal, $\phi_T^s \approx \phi_P^s$, and the direct CP asymmetry is proportional to $\sin(\phi_T^s - \phi_P^s)$ [23]. Besides, the time-integrated CP asymmetries in Table 2 differ dramatically from the corresponding direct CP asymmetries $-C_f$'s. According to Eq. (6), the differences mainly come from the large mixing parameter ΔM.

Table 2
Branching ratios (in units of 10^{-7}) and CP observables for the $\Delta S = 1$ $B_s^0 \to PT$ decays, whose final states are CP eigenstates.

Modes	Amplitudes	Br	C_f	D_f	S_f	time-inte A_{CP} (%)
$\pi^0 a_2^0$	PA	$0.90^{+0.19+0.31}_{-0.14-0.31}$	$-0.082^{+0.072+0.055}_{-0.001-0.015}$	$-0.988^{+0.003+0.001}_{-0.003-0.003}$	$-0.133^{+0.021+0.008}_{-0.031-0.011}$	$0.50^{+0.10+0.03}_{-0.10-0.03}$
$\pi^0 f_2$	A, PC	$0.048^{+0.012+0.002}_{-0.016-0.012}$	$-0.04^{+0.06+0.02}_{-0.12-0.06}$	$-0.66^{+0.08+0.08}_{-0.02-0.04}$	$0.75^{+0.06+0.06}_{-0.01-0.04}$	$-2.7^{+0.1+0.2}_{-0.2-0.2}$
$\pi^0 f_2'$	PC	$1.2^{+0.6+0.1}_{-0.5-0.1}$	$-0.05^{+0.01+0.01}_{-0.02-0.02}$	$-0.95^{+0.01+0.03}_{-0.01-0.02}$	$0.30^{+0.03+0.07}_{-0.02-0.07}$	$-1.0^{+0.1+0.3}_{-0.1-0.3}$
ηa_2^0	C, A	$0.047^{+0.013+0.010}_{-0.010-0.012}$	$0.02^{+0.01+0.01}_{-0.02-0.06}$	$0.40^{+0.01+0.06}_{-0.01-0.04}$	$0.92^{+0.01+0.02}_{-0.01-0.03}$	$-3.6^{+0+0.1}_{-0.01-0.1}$
ηf_2	PC	$9.8^{+2.7+3.2}_{-2.2-2.6}$	$-0.014^{+0.003+0.008}_{-0.008-0.010}$	$-0.995^{+0.001+0.002}_{-0.001-0}$	$-0.098^{+0.007+0.004}_{-0.007-0.020}$	$0.30^{+0.02+0.07}_{-0.02-0.01}$
$\eta f_2'$	PA	96^{+20+36}_{-19-30}	$0.022^{+0.004+0.003}_{-0.004-0.003}$	-1.000^{+0+0}_{-0-0}	$0.024^{+0.004+0.003}_{-0.004-0.005}$	$-0.10^{+0.02+0.02}_{-0.01-0.01}$
$\eta' a_2^0$	C, A	$0.13^{+0.03+0.03}_{-0.03-0.03}$	$0.03^{+0.01+0.02}_{-0.01-0.01}$	$0.28^{+0.03+0.04}_{-0-0.03}$	$0.96^{+0+0.01}_{-0.01-0.01}$	$-3.7^{+0+0.1}_{-0-0}$
$\eta' f_2$	PC	30^{+7+11}_{-7-10}	$-0.005^{+0+0.002}_{-0.012-0.010}$	$-0.994^{+0.001+0.001}_{-0.001-0.001}$	$-0.104^{+0.011+0.006}_{-0.006-0.006}$	$0.40^{+0.02+0.02}_{-0.04-0.02}$
$\eta' f_2'$	PA, PT	245^{+69+99}_{-59-84}	$-0.007^{+0.004+0}_{-0.003-0.001}$	-1.000^{+0+0}_{-0-0}	$-0.009^{+0.006+0.004}_{-0.002-0.001}$	$0.030^{+0.010+0.002}_{-0.020-0.010}$

Table 3
Branching ratios (in units of 10^{-7}) and CP observables for the rest $\Delta S = 1$ decays.

Modes	C_f	D_f	S_f	$C_{\bar f}$	$D_{\bar f}$	$S_{\bar f}$	Br	A_{CP} (%)
$\pi^+ a_2^-$	$-0.15^{+0.01+0.02}_{-0.04-0.05}$	$-0.98^{+0+0.01}_{-0-0.01}$	$-0.10^{+0.07+0.05}_{-0.01-0.01}$	$-0.05^{+0.07+0.07}_{-0.02-0.01}$	$-0.98^{+0.01+0.01}_{-0.01-0.01}$	$0.18^{+0.04+0.04}_{-0.02-0.03}$	$1.8^{+0.4+0.6}_{-0.2-0.8}$	13^{+3+5}_{-5-5}
$K^+ K_2^{*-}$	$0.49^{+0.07+0.02}_{-0.06-0.01}$	$-0.85^{+0.04+0}_{-0.03-0}$	$-0.18^{+0.02+0.03}_{-0.04-0.05}$	$0.03^{+0.11+0.09}_{-0.08-0.13}$	$-0.71^{+0.09+0.03}_{-0.06-0.02}$	$-0.70^{+0.07+0.03}_{-0.07-0.03}$	86^{+20+28}_{-16-24}	-28^{+2+5}_{-3-6}
$K^0 \bar{K}_2^{*0}$	$0.24^{+0.08+0.03}_{-0.06-0.05}$	$-0.91^{+0.03+0.02}_{-0.02-0.02}$	$-0.34^{+0.03+0.04}_{-0.03-0.03}$	$0.24^{+0.08+0.03}_{-0.06-0.05}$	$-0.91^{+0.03+0.02}_{-0.02-0.02}$	$-0.34^{+0.03+0.04}_{-0.03-0.03}$	70^{+14+24}_{-12-20}	0

There exist more complicated $\Delta S = 1$ modes, in which either a B_s^0 or $\bar{B}_s^0$ meson can decay into f and $\bar{f}$ with $f \neq \bar{f}$. Even though a final state is identified in this case, there is no way to determine whether the initial state is a B_s^0 or $\bar{B}_s^0$ meson directly. It is then difficult to distinguish the four channels in Eq. (3), and time-dependent measurements are also required. For experimental access, we define the CP asymmetry parameter only by charge-tag of final states

$$A_{CP} \equiv \frac{Br(B_s^0/\bar{B}_s^0(\infty) \to \bar{f}) - Br(B_s^0/\bar{B}_s^0(\infty) \to f)}{Br(B_s^0/\bar{B}_s^0(\infty) \to \bar{f}) + Br(B_s^0/\bar{B}_s^0(\infty) \to f)}. \tag{7}$$

All the CP observables, and the sum of the branching ratios of a pair of channels defined by

$$Br \equiv \frac{1}{2}\big[Br\big(B_s^0(\infty) \to f\big) + Br\big(\bar{B}_s^0(\infty) \to \bar{f}\big) + Br\big(B_s^0(\infty) \to \bar{f}\big) + Br\big(\bar{B}_s^0(\infty) \to f\big)\big], \tag{8}$$

are presented in Table 3. For the $B_s^0 \to \bar{K}^0 K_2^{*0}$ set, all the f-related CP observables are equal to the $\bar{f}$-related ones, and the CP asymmetry parameter A_{CP} is exactly zero. There are no tree contributions, and the penguin amplitudes share one common weak phase in these decays. It is then straightforward to arrive at $\lambda_f = \bar{\lambda}_{\bar{f}}$, and thus $C(D,S)_f = C(D,S)_{\bar{f}}$ and $A_{CP} = 0$.

In this Letter we have investigated the $B_s \to PT$ decays in the PQCD approach, whose branching ratios and CP asymmetry parameters were predicted. It was noticed that the absence of the factorizable tensor-emission amplitudes in these decays leads to differences from the $B_s \to PV$ ones. Owing to the significant B_s^0–$\bar{B}_s^0$ mixing effect, the time-integrated CP asymmetries have been redefined and calculated for the $\Delta S = 1$ modes. The U-spin symmetry was found to be considerably broken, when the $B_s^0 \to \pi^+ K_2^{*-}$ and $K^- a_2^+$ branching ratios are compared to the corresponding $B^0 \to K^+ a_2^-$ and $\pi^- K_2^{*+}$ ones. The branching ratios of some modes reach $\mathcal{O}(10^{-6})$ or even $\mathcal{O}(10^{-5})$, including $B_s^0 \to \eta f_2'$, $\eta' f_2$, $\eta' f_2'$, $K^+ K_2^{*-}$, $K^0 \bar{K}_2^{*0}$, and $\pi^+ K_2^{*-}$, which are expected to be measured at LHCb experiments. There is also potential to observe CP violation effects in the $B_s^0 \to \pi^+ K_2^{*-}$, $K^+ K_2^{*-}$ and $K^0 \bar{K}_2^{*0}$ decays in the near future.

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